

Overview of \bar{K} -Nuclear Theory and Phenomenology^{*)}

— Search for Narrow Quasibound States —

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Experimental evidence for \bar{K} -nuclear quasibound states is briefly reviewed. Theoretical and phenomenological arguments for and against deep \bar{K} -nucleus potentials which might allow for narrow quasibound states are reviewed, with recent calculations suggesting that $\Gamma_{\bar{K}} \geq 100$ MeV for $B_{\bar{K}} \leq 100$ MeV. Results of RMF calculations that provide a lower limit of $\Gamma_{\bar{K}} \sim 50 \pm 10$ MeV on the width of deeply bound states are discussed.

§1. Introduction

The \bar{K} -nucleus interaction near threshold is strongly attractive and absorptive as suggested by fits to the strong-interaction shifts and widths of K^- -atom levels.¹⁾ Global fits yield ‘deep’ density dependent (DD) optical potentials with nuclear-matter depth $\text{Re } V_{\bar{K}}(\rho_0) \sim -(150 - 200)$ MeV,^{2)–6)} whereas in the other extreme case several studies that fit the low-energy K^-p reaction data, including the $I = 0$ quasibound state $A(1405)$ as input for constructing self consistently DD optical potentials, obtain relatively ‘shallow’ potentials with $\text{Re } V_{\bar{K}}(\rho_0) \sim -(40 - 60)$ MeV.^{7)–9)} The issue of depth of $\text{Re } V_{\bar{K}}$ is discussed in Section 2 and the implications of a ‘deep’ potential for the existence and properties of \bar{K} -nucleus quasibound states are discussed in Section 3. Below we briefly discuss deeply bound K^- atomic states and review the debate over deeply bound \bar{K} nuclear states in light nuclei.

1.1. Deeply bound K^- atomic states

Paradoxically, due to the strong (absorptive) $\text{Im } V_{\bar{K}}$, relatively narrow K^- deeply bound *atomic* states are expected to exist,¹⁰⁾ independently of the size of $\text{Re } V_{\bar{K}}$. Figure 1 from Ref.¹¹⁾ shows on the left-hand side a calculated spectrum of K^- atomic states in ^{208}Pb where, in particular, all the circular states below the $7i$ ($l = 6$) state are not populated by X-ray transitions due to the strong K^- -nuclear absorption, and on the right-hand side it demonstrates saturation of the $2p$ atomic-state width for realistically large values of the absorptivity parameter $\text{Im } b_0$ of $V_{\bar{K}}$. The physics behind is that a strong $\text{Im } V_{\bar{K}}$ acts repulsively, suppressing the *atomic* wavefunction $\Psi_{\text{atom}}(\mathbf{r})$ in the region of overlap with $\text{Im } V_{\bar{K}}$, which according to the expression

$$\frac{\Gamma_{\text{atom}}}{2} = \frac{\int \text{Im}(-V_{\bar{K}}(r)) |\Psi_{\text{atom}}(\mathbf{r})|^2 d\mathbf{r}}{\int |\Psi_{\text{atom}}(\mathbf{r})|^2 d\mathbf{r}} \quad (1.1)$$

implies suppression of the width Γ_{atom} as well. The small calculated width of ‘deeply bound’ atomic states, $\Gamma_{\text{atom}} \leq 2$ MeV, also confirmed by the calculation of Ref.,¹²⁾ calls for experimental ingenuity to form these levels nonradiatively.¹³⁾

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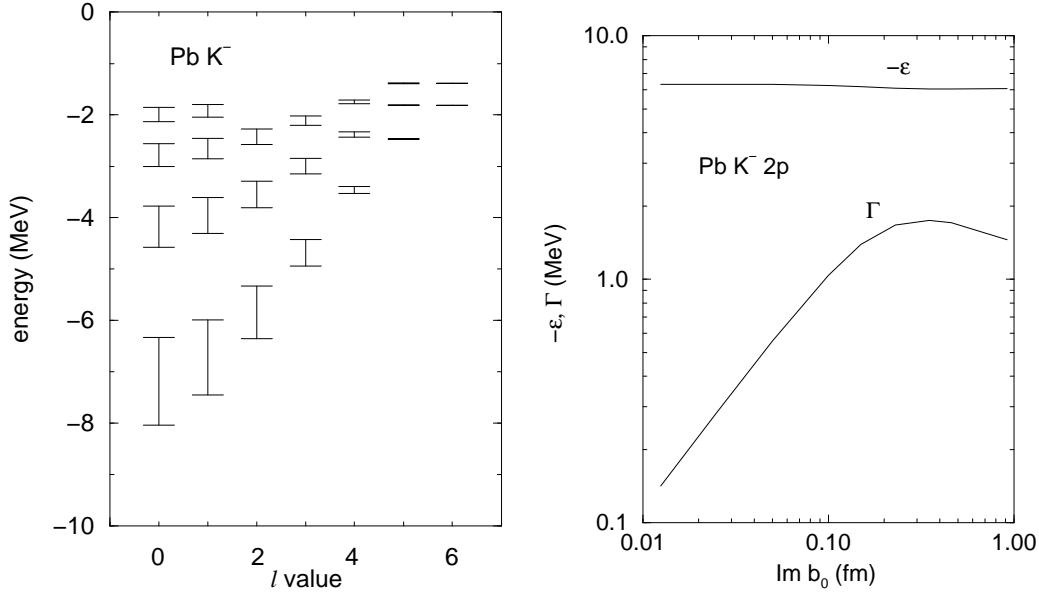


Fig. 1. Calculated K^- ‘deeply bound’ atomic states in ^{208}Pb (left) and saturation of width Γ_{atom} for the $2p$ ‘deeply bound’ state (right) as function of absorptivity, $\text{Im } b_0$, for $\text{Re } b_0 = 0.62$ fm.

1.2. Deeply bound K^- nuclear states in light nuclei

No saturation of $\Gamma_{\bar{K}}$ holds for \bar{K} -nuclear states which retain substantial overlap with the potential. In addition to asking (i) whether it is possible at all to bind *strongly* \bar{K} mesons in nuclei, one should ask (ii) are such quasibound states sufficiently narrow to allow observation and identification? The first question was answered affirmatively by Nogami¹⁴⁾ as early as 1963, estimating that the lightest system K^-pp is bound by about 10 MeV in its $I = 1/2$ state. The first calculation, by Yamazaki and Akaishi¹⁵⁾ gave binding energy $B \sim 50$ MeV and width $\Gamma \sim 60$ MeV. Preliminary AMD calculations by Doté and Weise^{16), 17)} which implicitly account for $\bar{K}N - \pi Y$ coupling obtain somewhat lower values of B with an estimated width $\Gamma \sim 100$ MeV. Coupled-channel $\bar{K}NN - \pi\Sigma N$ Faddeev calculations^{18), 19)} of K^-pp have confirmed this order of magnitude of binding, $B \sim 55 - 75$ MeV, differing on the width; the calculations of Ref.¹⁸⁾ give large values $\Gamma \sim 100$ MeV for the mesonic width. These Faddeev calculations overlook the $\bar{K}NN \rightarrow YN$ coupling to nonmesonic channels which are estimated to add, conservatively, 20 MeV to the overall width. Altogether, the widths calculated for the K^-pp quasibound state are likely to be as large as to make it difficult to identify it experimentally.

The current experimental and theoretical interest in \bar{K} -nuclear bound states was triggered back in 1999 by the suggestion of Kishimoto²⁰⁾ to look for such states in the nuclear reaction (K^-, p) in flight, and by Akaishi and Yamazaki^{21), 22)} who suggested to look for a $\bar{K}NNN$ $I = 0$ state bound by over 100 MeV for which the main $\bar{K}N \rightarrow \pi\Sigma$ decay channel would be kinematically closed.^{*)} Some controversial evidence for relatively narrow states was presented initially in (K_{stop}^-, n) and (K_{stop}^-, p)

^{*)} Wycech had conjectured that the width of such states could be as small as 20 MeV.²³⁾

reactions on ^4He (KEK-PS E471)^{(24),(25)} but has just been withdrawn (KEK-PS E549/570).⁽²⁶⁾ \bar{K} -nuclear states were also invoked to explain few statistically-weak irregularities in the neutron spectrum of the (K^-, n) in-flight reaction on ^{16}O (BNL-AGS, parasite E930⁽²⁷⁾), but subsequent (K^-, n) and (K^-, p) reactions on ^{12}C at $p_{\text{lab}} = 1 \text{ GeV}/c$ (KEK-PS E548⁽²⁸⁾) have not disclosed any peaks beyond the appreciable strength observed below the \bar{K} -nucleus threshold. Ongoing experiments by the FINUDA spectrometer collaboration at DAΦNE, Frascati, already claimed evidence for a relatively broad K^-pp deeply bound state ($B \sim 115 \text{ MeV}$) in K^-_{stop} reactions on Li and ^{12}C , by observing back-to-back Ap pairs from the decay $K^-pp \rightarrow Ap$,⁽²⁹⁾ but these pairs could more naturally arise from conventional absorption processes at rest when final-state interaction is taken into account.⁽³⁰⁾ Indeed, the $K^-_{\text{stop}}pn \rightarrow \Sigma^-p$ reaction on ^6Li has also been recently observed.⁽³¹⁾ Another recent claim of a very deep and narrow K^-pp state ($B \sim 160 \text{ MeV}$, $\Gamma \sim 30 \text{ MeV}$) is also based on observing decay Ap pairs, using \bar{p} annihilation data on ^4He from the OBELIX spectrometer at LEAR, CERN.⁽³²⁾ The large value of B_{K^-pp} over 100 MeV conjectured by these experiments is at odds with *all* the few-body calculations of the K^-pp system listed above. Could Ap pairs assigned in the above analyses to K^-pp decay in fact result from nonmesonic decays of different clusters, say the $\bar{K}NNN$ $I = 0$ quasibound state? A definitive identification of the $\{\bar{K}[NN]_{I=1}\}_{I=1/2}$ quasibound state may optimally be reached through fully exclusive formation reactions, such as:

$$K^- + ^3\text{He} \rightarrow n + \{\bar{K}[NN]_{I=1}\}_{I=1/2, I_z=+1/2}, \quad p + \{\bar{K}[NN]_{I=1}\}_{I=1/2, I_z=-1/2}, \quad (1\cdot2)$$

the first of which is scheduled for day-one experiment in J-PARC.⁽³³⁾ Finally, preliminary evidence for a $\bar{K}NNN$ $I = 0$ state with $B = 67 \pm 5 \text{ MeV}$, $\Gamma = 37 \pm 10 \text{ MeV}$ has been recently presented by the FINUDA collaboration on ^6Li by observing back-to-back Ad pairs.⁽³⁴⁾ It is clear that the issue of \bar{K} nuclear states is far yet from being experimentally resolved and more dedicated, systematic searches are necessary.

§2. \bar{K} -nucleus potentials

2.1. Chirally motivated models

The Born approximation for $V_{\bar{K}}$ due to the leading-order Tomozawa-Weinberg (TW) vector term of the chiral effective Lagrangian⁽³⁵⁾ yields a sizable attraction:

$$V_{\bar{K}} = -\frac{3}{8f_\pi^2} \rho \sim -55 \frac{\rho}{\rho_0} \text{ MeV} \quad (2\cdot1)$$

for $\rho_0 = 0.16 \text{ fm}^{-3}$, where $f_\pi \sim 93 \text{ MeV}$ is the pseudoscalar meson decay constant. Iterating the TW term plus next-to-leading-order terms, within an *in-medium* coupled-channel approach constrained by the $\bar{K}N - \pi\Sigma - \pi\Lambda$ data near the $\bar{K}N$ threshold, roughly doubles this \bar{K} -nucleus attraction. It is found (e.g. Ref.⁽³⁶⁾) that the $\Lambda(1405)$ quickly dissolves in the nuclear medium at low density, so that the repulsive free-space scattering length a_{K^-p} , as function of ρ , becomes *attractive* well below ρ_0 . Since the attractive $I = 1$ a_{K^-n} is only weakly density dependent, the

resulting in-medium $\bar{K}N$ isoscalar scattering length $b_0(\rho) = \frac{1}{2}(a_{K-p}(\rho) + a_{K-n}(\rho))$ translates into a strongly attractive $V_{\bar{K}}$:

$$V_{\bar{K}}(r) = -\frac{2\pi}{\mu_{KN}} b_0(\rho) \rho(r), \quad \text{Re } V_{\bar{K}}(\rho_0) \sim -110 \text{ MeV}. \quad (2.2)$$

However, when $V_{\bar{K}}$ is calculated *self consistently*, namely by including $V_{\bar{K}}$ in the propagator G_0 used in the Lippmann-Schwinger equation determining $b_0(\rho)$, one obtains $\text{Re } V_{\bar{K}}(\rho_0) \sim -(40 - 60) \text{ MeV}$.^{(7)–(9)} The main reason for this weakening of $V_{\bar{K}}$, approximately back to that of Eq. (2.1), is the strong absorptive effect which $V_{\bar{K}}$ exerts within G_0 to suppress the higher Born terms of the $\bar{K}N$ TW potential.

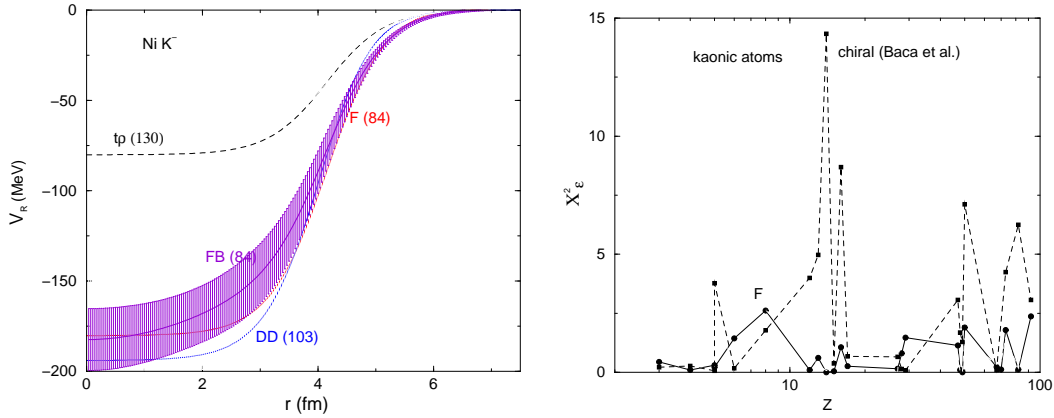


Fig. 2. Real part of the $\bar{K}^- - ^{58}\text{Ni}$ potential obtained in a global fit to K^- -atom data using the model-independent FB technique⁽⁶⁾, in comparison with other best-fit potentials and χ^2 values in parentheses (left), and contributions to the χ^2 of K^- atomic shifts for the *deep* potential F from Ref.⁽⁵⁾ and for the *shallow* chirally-based potential from Ref.⁽¹²⁾ (right).

2.2. Fits to K^- -atom data

The K^- -atom data used in global fits⁽¹⁾ span a range of nuclei from ^7Li to ^{238}U , with 65 level-shift, -width and -yield data points. Figure 2 shows on the left-hand side fitted $\text{Re } V_{\bar{K}}$ for ^{58}Ni , for a $t\rho$ potential with a complex strength t , for two density-dependent potentials⁽⁵⁾ marked by DD and F, and for a *model independent* potential obtained by adding a Fourier-Bessel (FB) series to a $t\rho$ potential.⁽⁶⁾ The depth of the $t\rho$ potential is about 70 – 80 MeV, whereas the density-dependent potentials (including FB) are considerably deeper, 150 – 200 MeV. Although DD and F have very different parameterizations, the resulting potentials are quite similar to each other. These latter potentials yield substantially lower χ^2 values (shown in the figure) than the χ^2 value for the $t\rho$ potential. The superiority of the *deep* potential F to a *shallow* potential based on the self-consistent chiral model of Ramos and Oset⁽⁸⁾ in reproducing the K^- -atom level shifts is demonstrated on the right-hand side of Fig. 2. In particular, the shape of potential F departs appreciably from $\rho(r)$ for $\rho(r)/\rho_0 \leq 0.2$, where the physics of the $\Lambda(1405)$ plays a role. It is a challenge for chiral-model practitioners to match the quality of potential F's atomic fit.

2.3. Further considerations

(i) QCD sum-rule estimates³⁷⁾ for vector (v) and scalar (s) self-energies give:

$$\Sigma_v(\bar{K}) \sim -\frac{1}{2} \Sigma_v(N) \sim -\frac{1}{2} (200) \text{ MeV} = -100 \text{ MeV}, \quad (2.3)$$

$$\Sigma_s(\bar{K}) \sim \frac{m_s}{M_N} \Sigma_s(N) \sim \frac{1}{10} (-300) \text{ MeV} = -30 \text{ MeV}, \quad (2.4)$$

where m_s is the strange-quark (current) mass. The factor 1/2 in Eq. (2.3) is due to the one nonstrange antiquark \bar{q} in the \bar{K} out of two possible, and the minus sign is due to G -parity going from q to \bar{q} . This rough estimate gives then $V_{\bar{K}}(\rho_0) \sim -130 \text{ MeV}$. The QCD sum-rule approach essentially refines the mean-field argument^{38), 39)}

$$V_{\bar{K}}(\rho_0) \sim \frac{1}{3} (\Sigma_s(N) - \Sigma_v(N)) \sim -170 \text{ MeV}, \quad (2.5)$$

where the 1/3 factor is again due to the one nonstrange antiquark in the \bar{K} , but here with respect to the three nonstrange quarks of the nucleon.

(ii) The ratio of K^-/K^+ production cross sections in nucleus-nucleus and proton-nucleus collisions near threshold, measured by the Kaon spectrometer (KaoS) collaboration at SIS, GSI, gives clear evidence for a strongly attractive $V_{\bar{K}}$, estimated⁴⁰⁾ as $V_{\bar{K}}(\rho_0) \sim -80 \text{ MeV}$ by relying on BUU transport calculations normalized to the value $V_K(\rho_0) \sim +25 \text{ MeV}$. Since $\bar{K}NN \rightarrow YN$ absorption apparently was disregarded in these calculations, a deeper $V_{\bar{K}}$ may follow when it is included.

§3. RMF dynamical calculations

3.1. \bar{K} -nucleus RMF methodology

In this model, expanded in Ref.⁵⁾ the (anti)kaon interaction with the nuclear medium is incorporated by adding to \mathcal{L}_N the Lagrangian density \mathcal{L}_K :

$$\mathcal{L}_K = \mathcal{D}_\mu^* \bar{K} \mathcal{D}^\mu K - m_K^2 \bar{K} K - g_{\sigma K} m_K \sigma \bar{K} K. \quad (3.1)$$

The covariant derivative $\mathcal{D}_\mu = \partial_\mu + ig_{\omega K} \omega_\mu$ describes the coupling of the (anti)kaon to the vector meson ω . The (anti)kaon coupling to the isovector ρ meson was found to have negligible effects. The \bar{K} meson induces additional source terms in the equations of motion for the meson fields σ and ω_0 . It thus affects the scalar $S = g_{\sigma N} \sigma$ and the vector $V = g_{\omega N} \omega_0$ potentials which enter the Dirac equation for nucleons, and this leads to rearrangement or polarization of the nuclear core, as shown on the left-hand side of Fig. 3 for the calculated average nuclear density $\bar{\rho} = \frac{1}{A} \int \rho^2 d\mathbf{r}$ as a function of B_{K^-} for K^- nuclear $1s$ states across the periodic table. It is seen that in the light K^- nuclei, $\bar{\rho}$ increases substantially with B_{K^-} to values about 50% higher than without the \bar{K} .*) Furthermore, in the Klein-Gordon equation satisfied by the \bar{K} , the scalar $S = g_{\sigma K} \sigma$ and the vector $V = -g_{\omega K} \omega_0$ potentials become *state*

*) The increase of the central nuclear densities is bigger, up to 100%, and is nonnegligible even in the heavier K^- nuclei where it is confined to a small region of order 1 fm.

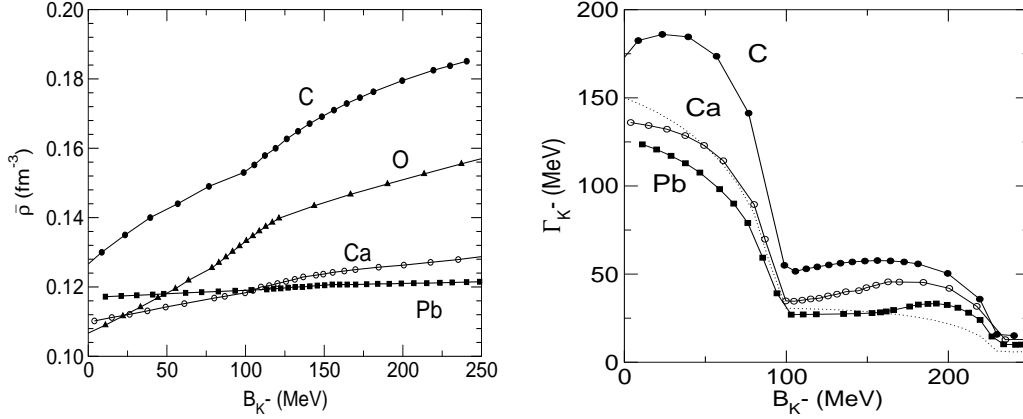


Fig. 3. Dynamically calculated average density $\bar{\rho}$ (left) and widths Γ_{K^-} (right) of $1s$ K^- -nuclear states in the nuclei denoted, as function of the $1s$ K^- binding energy.⁵⁾

dependent through the *dynamical* density dependence of the mean-field potentials S and V , as expected in a RMF calculation. An imaginary $\text{Im } V_{\bar{K}} \sim t\rho$ was added, fitted to the K^- atomic data.⁴⁾ It was then suppressed by an energy-dependent factor $f(B_{\bar{K}})$, considering the reduced phase-space for the initial decaying state and assuming two-body final-state kinematics for the decay products in the $\bar{K}N \rightarrow \pi Y$ mesonic modes (80%) and in the $\bar{K}NN \rightarrow YN$ nonmesonic modes (20%).

The RMF coupled equations were solved self-consistently. For a rough idea, whereas the static calculation gave $B_{K^-}^{1s} = 132$ MeV for the K^- $1s$ state in ^{12}C , using the values $g_{\omega K}^{\text{atom}}$, $g_{\sigma K}^{\text{atom}}$ corresponding to the K^- -atom fit, the dynamical calculation gave $B_{K^-}^{1s} = 172$ MeV. In order to scan a range of values for $B_{K^-}^{1s}$, $g_{\sigma K}$ and $g_{\omega K}$ were varied in given intervals of physical interest.

3.2. Binding energies and widths

Beginning approximately with ^{12}C , the following conclusions may be drawn: (i) For given values of $g_{\sigma K}$, $g_{\omega K}$, the \bar{K} binding energy $B_{\bar{K}}$ saturates, except for a small increase due to the Coulomb energy (for K^-). (ii) The difference between the binding energies calculated dynamically and statically, $B_{\bar{K}}^{\text{dyn}} - B_{\bar{K}}^{\text{stat}}$, is substantial in light nuclei, increasing with $B_{\bar{K}}$ for a given value of A , and decreasing monotonically with A for a given value of $B_{\bar{K}}$. It may be neglected only for very heavy nuclei. The same holds for the nuclear rearrangement energy $B_{\bar{K}}^{\text{s.p.}} - B_{\bar{K}}$ which is a fraction of $B_{\bar{K}}^{\text{dyn}} - B_{\bar{K}}^{\text{stat}}$. (iii) The width $\Gamma_{\bar{K}}(B_{\bar{K}})$ decreases monotonically with A , as shown in the right-hand side of Fig. 3 for $1s$ states. The dotted line shows the static 'nuclear-matter' limit corresponding to the K^- -atom fitted value $\text{Im } b_0 = 0.62$ fm and for $\rho(r) = \rho_0 = 0.16$ fm^{-3} , using the same phase-space suppression factor as in the 'dynamical' calculations. It is clearly seen that the functional dependence $\Gamma_{K^-}(B_{K^-})$ follows the shape of the dotted line. This dependence is due primarily to the binding-energy dependence of $f(B_{K^-})$ which falls off rapidly until $B_{K^-} \sim 100$ MeV, where the dominant $\bar{K}N \rightarrow \pi\Sigma$ gets switched off, and then stays rather flat in the range $B_{K^-} \sim 100$ – 200 MeV where the width is dominated by the $\bar{K}NN \rightarrow YN$ absorption

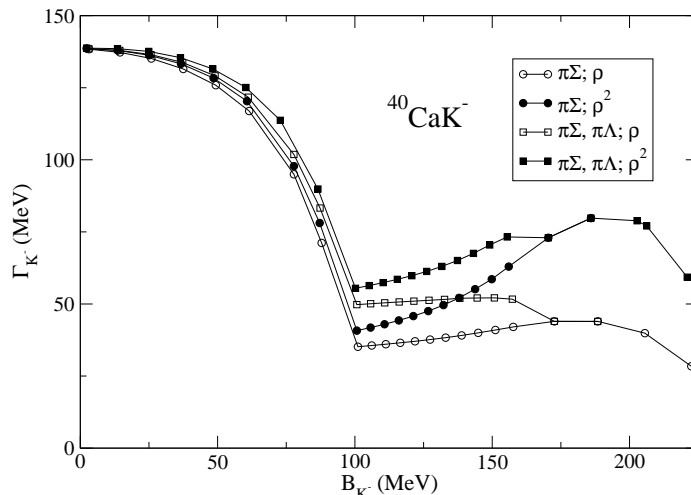


Fig. 4. Dynamically calculated widths of the $1s$ \bar{K}^- -nuclear state in ^{40}Ca for $\pi\Sigma + \pi\Lambda$ vs. $\pi\Sigma$ final mesonic absorption channels, and for ρ^2 vs. ρ dependence of the final nonmesonic absorption channels, as function of the \bar{K}^- binding energy, from Gazda et al.⁵⁾

modes. The widths calculated in this range are considerably larger than given by the dotted line (except for Pb in the range $B_{K^-} \sim 100 - 150$ MeV) due to the dynamical nature of the RMF calculation, whereby the nuclear density is increased by the polarization effect of the \bar{K}^- . Adding perturbatively the residual width neglected in this calculation, partly due to the $\bar{K}N \rightarrow \pi\Lambda$ secondary mesonic decay channel, a lower limit $\Gamma_{\bar{K}} \geq 50 \pm 10$ MeV is obtained for the binding energy range $B_{K^-} \sim 100 - 200$ MeV. Figure 4 shows the effect of splitting the 80% mesonic decay width, previously assigned to $\pi\Sigma$, between $\pi\Sigma$ (70%) and $\pi\Lambda$ (10%), and also of simulating the 20% nonmesonic absorption channels by a ρ^2 dependence compared to $\text{Im } V_{\bar{K}} \sim t\rho$ used by Mareš et al.⁵⁾ These added contributions make the above lower limit $\Gamma_{\bar{K}} \geq 50 \pm 10$ MeV a rather conservative one.

§4. Conclusions and acknowledgements

I have reviewed the phenomenological and theoretical evidence for a substantially attractive \bar{K} -nucleus interaction potential, from a ‘shallow’ potential of depth 40–60 MeV to a ‘deep’ potential of depth 150–200 MeV at nuclear-matter density. In particular, I demonstrated the considerable extent to which the *deep* potentials are supported by \bar{K}^- atomic data fits. I then reported on recent *dynamical* calculations⁵⁾ for deeply quasibound \bar{K} -nuclear states across the periodic table. Substantial polarization of the core nucleus was found in light nuclei, but the ‘high’ densities reached are considerably lower than those found in the few-body calculations due to Akaishi, Yamazaki and collaborators.^{15), 22), 41)} The width $\Gamma_{\bar{K}}$ of quasibound states, as function of the binding energy $B_{\bar{K}}$, reflects the decay phase-space suppression on top of the increase provided by the density of the compressed nuclear core, resulting in a lower limit $\Gamma_{\bar{K}} \geq 50 \pm 10$ MeV, particularly useful for $B_{\bar{K}} \geq 100$ MeV.

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